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Electromagnetic Transient Fields from Spectrally Finite Exoatmospheric Photon Bursts

by
Arnold D. Levine
Department of Physics, West Virginia University

Andre H. Rausch
National Communications System

Ronald J. Chase
Army Research Laboratory

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Arnold D. Levine (West Virginia University), Andre H. Rausch (National Communications System), and Ronald J. Chase (Army Research Laboratory)

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Attn: AMSRL-WT-ND
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The old problem of electromagnetic transients from exoatmospheric photon bursts has been restudied, by considering the photons to be generated from a spectrally finite source and not a monoenergetic source, as is usually done. Variables in the electromagnetic field equations are obtained by spectral averaging. Stratification of target molecules in the atmosphere leads to stratification of radiation sources. A conducting layer in the lowest strata develops due to high-energy electrons, leading to a reduction in field intensities near the ground.

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Basic models for electromagnetic transients arising in exoatmospheric photon bursts were developed 25 years ago [1,2]. Explicit numerical calculations have been carried out since then [3,4]. These calculations, based on a "nominal" (i.e., monoenergetic) photon source, predict rather large field intensities, even at the ground level. We find, however, that this result is not to be expected if we assume a spectrally finite photon source. In typical applications, a spectrally finite case is more likely [5, p 36]. One might presume that, even in such a case, the monochromatic source is the proper one, since kinetic processes involving the electromagnetic field sources (x-ray photoproduction, Compton scattering, electron positron pair production) are quantum processes and may therefore be taken to be incoherent. Nevertheless, the time scale for these processes ($<10^{-12}$ s) is much smaller than the transient pulse time scale ($\sim 10^{-8}$ s), so the fields may be assumed to originate from coherent sources.

The problem under consideration may be envisioned if we start with a wave equation for the transverse fields in the impulse approximation [6]:

$$\frac{\partial[r\mathbf{E}(\mathbf{r},\tau)]}{\partial r} = \frac{-Z_0 r\mathbf{j}(\mathbf{r},\tau)}{2} , \tag{1}$$

where τ is the retarded time $[t-(r/c)]$ and Z_0 is $\mu_0 c$, the impedance of free space. The photon pulse gives rise to a primary electron current, which depends on the kinetic process for electron production. In turn, the primary current generates a secondary current through the process of ionization by collision with ambient molecules. The magnitude of the secondary current, which is generally proportional to the primary current, will normally build up to the magnitude of the primary current, so that the field intensity from that point onward propagates as a free wave whose magnitude is essentially independent of the source. This point, then, is referred to as the saturation limit. We find this condition schematically by setting the ratio of the magnitude of the primary current to that of the secondary current equal to unity. If we use the expression $j_p = -en_p v_p$ for the primary current, and the expression $j_s = -en_s v_s$ for the secondary current, we have for the saturation condition:

$$\left(\frac{n_s}{n_p}\right)\left(\frac{v_s}{v_p}\right) = 1 . \tag{2}$$

The first term on the left is obtained from the well-known formula for ionization by collision [7]:

$$\frac{d}{d\xi}\left(\frac{dn_s}{dn_p}\right) = \Gamma(\beta) , \tag{3}$$

where ξ is a path length, and $\Gamma(\beta)$ is a function of $\beta = (v_p/c)$ and the other parameters given by Fermi [7]. We then have

$$\frac{n_s}{n_p} = \int_0^{t_0} \frac{\Gamma(\beta)\beta c d\tau}{1-\beta} \; . \tag{4}$$

We integrate over the time profile of the incident pulse (in retarded time), and use for τ a value equal to $t_0 = 10^{-8}$ s. Thus, equation (2) becomes

$$\left[\frac{\Gamma(\beta)t_0}{1-\beta}\right] v_s(E) = 1 \; , \tag{5}$$

and the order of magnitude of the saturated field is determined as soon as we establish the relation between the drift velocity and the field intensity.

Note that this result is essentially independent of the intensity of the source and that the role of the geomagnetic field is of secondary importance. Only in the case of x-ray photoelectrons do we have to consider the details of the synchrotron motion in the impulse limit. For relativistic electrons, the motion in the impulse limit is nearly a straight line, and the role of the geomagnetic field is less significant in the sense that it is the determining factor for the directional effects, but not the field intensity. Since the most intense fields are produced by the relativistic electrons, we will ignore the effect of synchrotron motion in our discussion.

The relation between drift velocity and field intensity needed in equation (5) is rather complicated. The simplest possibility that can be used as a basis for discussion is a linear mobility model:

$$v_s = \mu E \; , \tag{6}$$

so that, finally, from equation (5) we obtain for the saturation field

$$E = \frac{1-\beta}{\Gamma(\beta)\mu t_0} \; . \tag{7}$$

The value for mobility of 6 m^2/V-s is suggested by Longmire et al [4]. It is found to be surprisingly good for the conditions under consideration because mobility is expected to be both field- and time-dependent. The drift velocities as normally reported are obtained using quiescent conditions, and thus may not be suitable for use in transient calculations [8]. The procedure used in the numerical work reported by Longmire et al [4] was that of calculating the mobility at each stage, and updating its value in each iteration. Detailed numerical calculations so obtained agreed in order of magnitude with the numerical estimates found by using a fixed value for mobility. We will use the value of 6 m^2/V-s, as suggested, in making our own estimates, and justify this approximation from another viewpoint later on.

Our contention is that the procedure of using a monoenergetic photon source leads to an overestimate of the field intensities to be expected at low altitudes. Photons that propagate down toward the earth encounter an increasingly dense array of target molecules with molecular density $n(z)$ at

altitude z, determined by the Law of Atmospheres for the regions of interest [9],

$$n(z) = n(0)\, e^{-\zeta} , \tag{8}$$

where $\zeta = (z/\Lambda)$, with Λ a parameter whose value is about 7 km, and $n(0)$ is the molecular density at ground level. As photons propagate into the lower atmosphere, reaction rates increase and the photon flux decreases. Ultimately a peak reaction level is reached, where reaction rates are maximum [10]:

$$e^{\overline{\zeta(w)}} = \frac{\Lambda}{\lambda(w)} , \tag{9}$$

where $\lambda(w)$ is the mean free path for collisions at ground level, i.e., $\lambda(w) = [1/(n(0)\sigma(w))]$. (We omit directionally dependent effects in the above relation by restricting ourselves to observations of pulses that occur directly overhead.) The energy-dependent cross section is equal to the Thompson cross section, σ_T, multiplied by a parameter $G(w)$ with $w = (h\nu/mc^2)$, i.e., $\sigma(w) = \sigma_T G(w)$ [11], so that

$$e^{\overline{\zeta(w)}} = G(w) e^{\zeta_T} . \tag{10}$$

Here, ζ_T is obtained from equation (8) using the Thompson cross section. The numerical value of ζ_T obtained this way is 5.2. This corresponds to a height of around 36 km. $G(w)$ is a monotonically decreasing function of w. For x-ray photoproduction $G(w) > 1$, and the reaction zone ranges from an altitude of around 80 km to about 36 km. For Compton scattering $G(w) < 1$, and the reaction zone ranges from about 36 km to about 18 km. The high-intensity fields created at altitudes of 30 to 36 km are attenuated by the secondary currents created by the high-energy electrons at altitudes of 18 to 24 km. This can be seen by referring to equation (7) and noting the $\beta \approx 1$ for ultrarelativistic electrons that are produced at these altitudes.

We proceeded in our estimate of the field intensities by replacing the monoenergetic source with a distributed one characterized by a spectral density $n(w)$ such that $\int dw\, n(w) = 1$, with $n(w)$ determined from the data in figure 1 [5, p 361]. Current density in the reaction zone was calculated by making a saddle-point expansion of variables around the height defined in equation (9) and using spectral averages, i.e., $j_p = \int dw\, n(w) j_p(w)$ and $j_s = \int dw\, n(w) j_s(w)$, etc. We obtained primary current densities that were fairly uniform over the range of altitudes in the reaction zone, but the secondary currents in the reaction zone decrease rapidly with altitude. This leads to an attenuation of the field intensity as it propagates down to the Earth's surface.

Numerically computed saturation field levels directly below the source location (linear mobility model) are presented in figure 2. The calculations were made using a mobility of 6 m^2/V-s, and the results for the ground intensities were found to be severely attenuated by the secondary electrons associated with the high-energy tail of the photon spectrum, i.e., for photons above 8 MeV. Our numerical estimate for the ground-level field intensity is about 6300 V/m.

Figure 1. Spectrum of initial photon energy 2000 yd from 20-kT explosion [5, p 361].

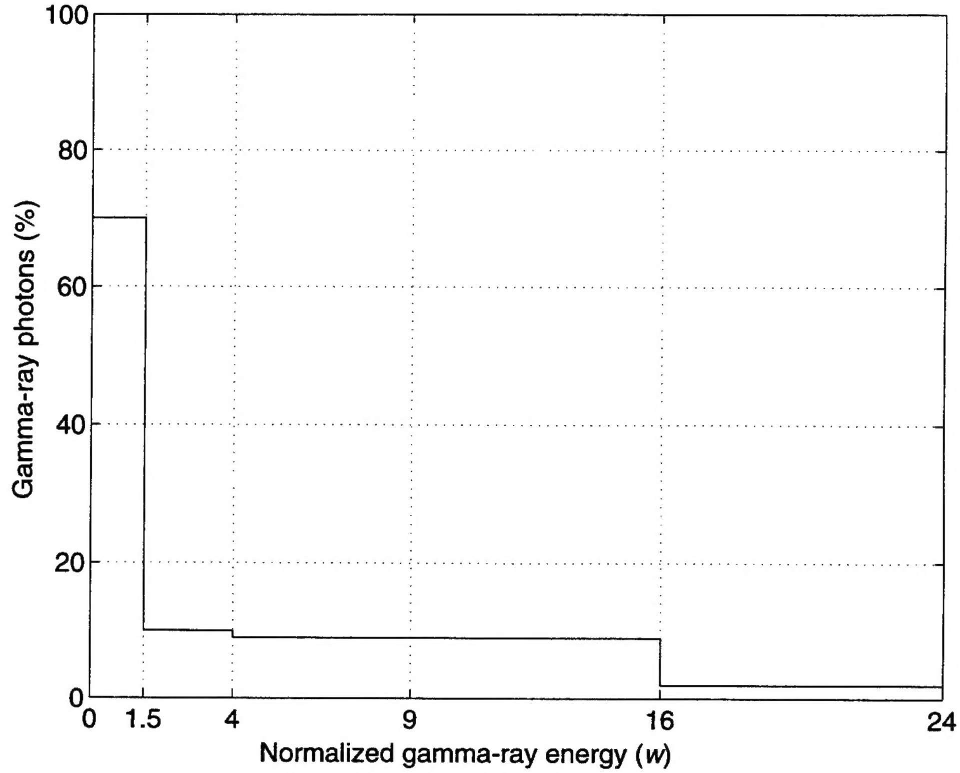


Figure 2. Saturated electric field versus photon energy for linear and nonlinear models. Field value for particular photon energy value represents magnitude of saturated electric field that would be found if photon spectral distribution were truncated for all larger photon energy values.

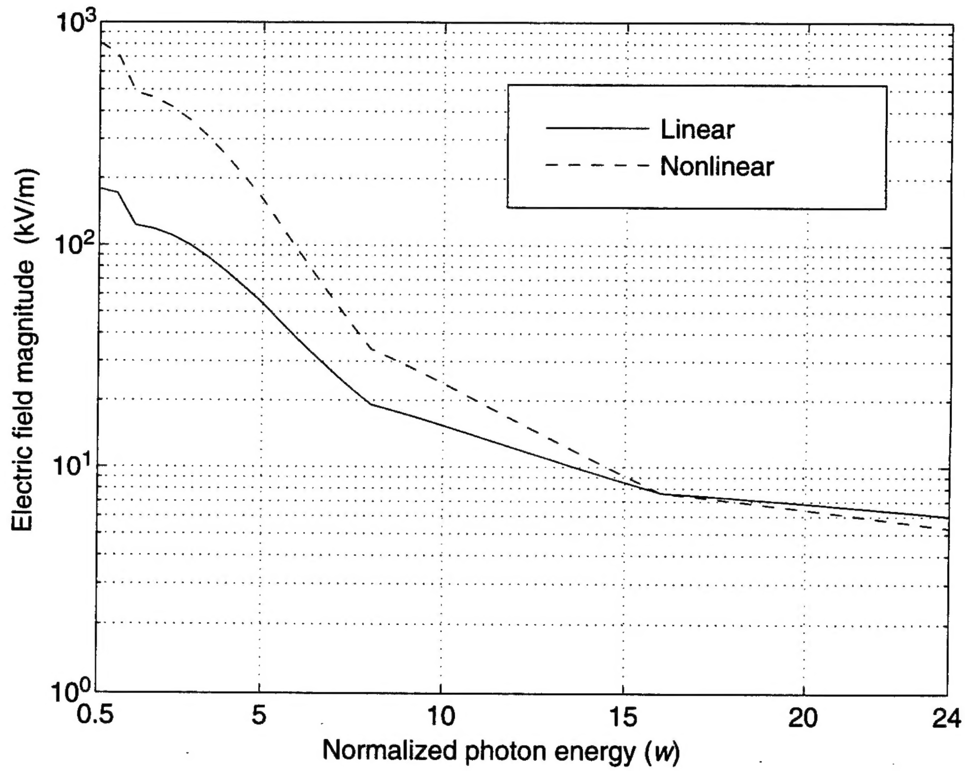
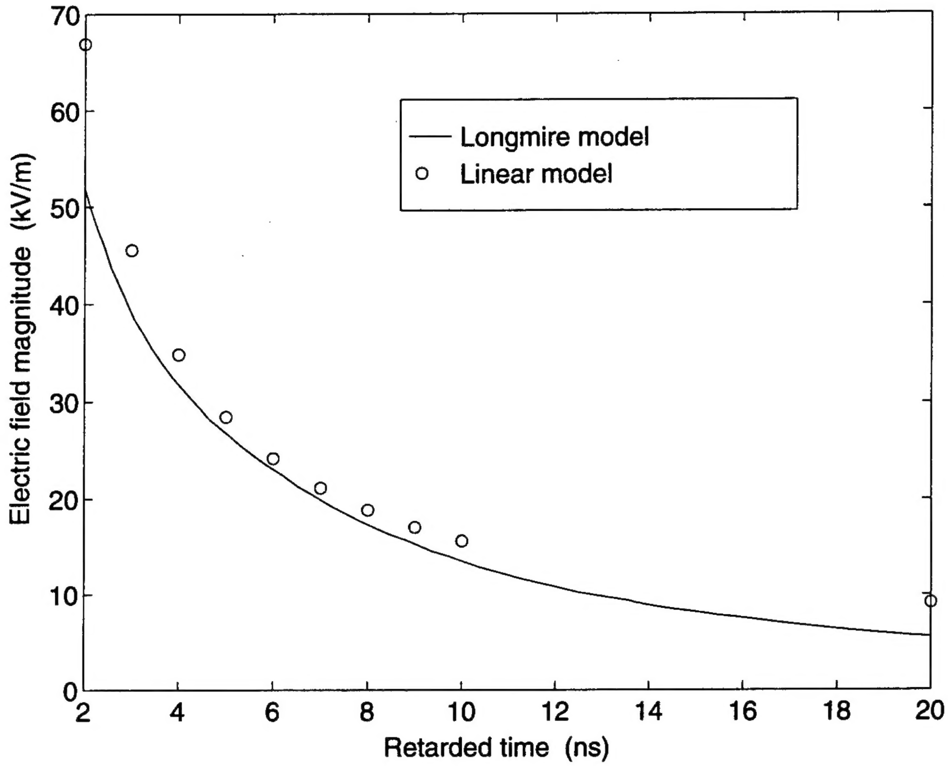


Figure 3. Comparison of saturated electric field predicted by our linear model and Longmire model for 1.6-MeV mono-energetic gamma source. Longmire data obtained from Longmire [3, fig. 13].



Apart from including the effects of the complete photon spectrum, our model does not differ significantly from the Longmire model [3]. By using a limiting procedure, our method recovers the monochromatic results for an impulse source of 1.6 MeV gamma energy. A comparison of the predicted saturation field level is presented in figure 3 for both models (data obtained from Longmire [3], fig. 13).

An improved estimate of the swarm parameters is possible if one considers the pseudo transport model,

$$\frac{dv}{dt} = -\frac{eE}{m} - \frac{v}{t_C} , \tag{11}$$

where t_C is a collision time. If, for t_C we take (λ/v), we find for the steady-state drift velocity the well-known result [12],

$$v = CE^{(1+\alpha)/2} , \tag{12}$$

where α represents some velocity dependence of the collision cross section. Since C is density-dependent, one finds

$$v = K\left(\frac{E}{p}\right)^{(1+\alpha)/2} , \tag{13}$$

with K a constant and p in atmospheres, to be a better representation. Ultimately one has, empirically, for oxygen, a good fit for the drift velocity, given by

$$v = 100\left(Ee^{\overline{\zeta(w)}}\right)^{1/2} . \tag{14}$$

From this we can calculate field mobilities and we find the value of 6 m^2/V-s to be a good approximation for the range of values in the region of interest. We can also use these parameters to obtain an estimate of the time for equilibrium to set in for the plasma. We find this to be less than 10^{-9} s, except for electrons of lowest energy. It is thus possible to go back to equation (1) and use equation (14) to calculate the saturation field (with, of course, spectrally averaged quantities). We can set

$$j_{\text{sec}} = -en_{\text{sec}}\chi\sqrt{E} \ , \tag{15}$$

with $\chi = 1400\sqrt{G(w)}$ in MKS units.

Since the wave equation is now nonlinear and not even separable, this is a major task. However, along geomagnetic flux lines, the equation is approximately separable. From equation (2), we develop a model wave equation of the form

$$\Xi \, d\Xi = -(\Xi - 1)\frac{dr}{r_\alpha} \ , \tag{16}$$

with $\Xi = \sqrt{E / E_1}$ and

$$\sqrt{E_1} = \frac{\beta_{\text{prim}} n_{\text{prim}} c}{n_{\text{sec}} \chi} \ ; \quad r_\alpha = \frac{4\sqrt{E_1}}{Z_0 e n_{\text{sec}} \chi} \ . \tag{17}$$

The solution to equation (16) is

$$e^{\Xi}(1 - \Xi) = \exp -\left(\frac{r_i - r}{r_\alpha}\right) . \tag{18}$$

We calculate r_α from equation (17) using spectral averages and we find it to be about 7.8 km. Typically $(r_i - r) \approx 20$ km, which makes the right-hand side $e^{-2.6}$ or ≈ 0.08, and $\Xi \approx 1$ in equation (16). By the use of spectral averages we find in equation (16) that $E_1 \approx 7000$ V/m.

One sees here that the saturation condition $\Xi \equiv 1$ is not fulfilled. As a matter of fact, Ξ ranges in value from about 0.5 to 1 over the range of parameters of interest. Nevertheless, figure 2 indicates that this model provides consistent results with the linear mobility model when the source includes a high-energy tail.

We have modified the standard model [3,4] for predicting the electric field intensity from an exoatmospheric gamma source by including the effects of all components of the photon pulse. We have found that the high-energy components of the photon pulse, which constitute only a small fraction of the total spectrum, are responsible for a reduction in the field at ground level.

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